

Maximally Natural Supersymmetry

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We consider 4D weak scale theories arising from 5D supersymmetric (SUSY) theories with maximal Scherk-Schwarz breaking at a Kaluza-Klein (KK) scale of several TeV. Many of the problems of conventional SUSY are avoided. Apart from 3rd family sfermions the SUSY spectrum is heavy, with only $\sim 50\%$ tuning at a gluino mass of ~ 2 TeV and a stop mass of ~ 650 GeV. A single Higgs doublet acquires a vacuum expectation value, so the physical Higgs is automatically Standard-Model-like. A new $U(1)'$ interaction raises m_h to 126 GeV. For minimal tuning the associated Z' , as well as the 3rd family sfermions, must be accessible to LHC13. A gravitational wave signal consistent with BICEP2 is possible if inflation occurs when the extra dimensions are small.

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The LHC has set stringent limits on the masses of SUSY particles and deviations in Higgs properties, implying a tuning of electroweak symmetry breaking (EWSB) at the percent level or worse for traditional SUSY models [1–6]. This undermines the motivation for SUSY as the solution to the hierarchy problem and the case for discovery of SUSY at the LHC or proposed future colliders. Given the importance of this issue for current and future searches for new physics we examine the possibility of constructing natural, untuned theories. Specifically, we consider 4D theories of the weak scale that arise from 5D SUSY theories with Scherk-Schwarz SUSY breaking (SSSB) at a KK scale $1/R$ of several TeV [7–20]. The key features are:

- The theory is never well approximated by a 4D softly-broken $N = 1$ SUSY limit. Many of the problems of the MSSM and its extensions are avoided.
- Higgsinos, gauginos, and the 1st and 2nd family sfermions get (mainly Dirac) SSSB masses of size $1/2R$.
- A natural SUSY spectrum [21–23] is obtained through localization of the 3rd family on a 4D brane. The absence of large logs due to the super-softness of SSSB [24–30] protects the weak scale and suppresses the tendency of the gluino to pull up the stop mass [2, 3].
- The μ term neither exists nor is needed. Only H_u acquires a VEV, and the down-like quark and lepton masses are generated by Kahler couplings to H_u^\dagger [31]. The physical Higgs is automatically SM-like.
- An additional SUSY breaking sector is necessarily present for radius stabilization with zero cosmological constant (CC), and SUSY breaking in this sector can naturally be driven by SSSB. Higher dimensional couplings of the MSSM fields to this sector play a crucial role in EWSB and collider phenomenology.
- A $U(1)'$ broken in this additional sector raises the Higgs mass to 126 GeV through an unusual non-decoupling D-term, with a Z' mass of order $1/R$.

The pattern of localization of matter and Higgs multiplets and the mechanism driving EWSB, generating

Yukawa couplings, and accommodating the observed physical Higgs mass lead to important differences from previously studied models of SSSB [7–20].

NATURAL SPECTRUM FROM SCHERK-SCHWARZ

Symmetries may be broken in a way preserving 4D Poincare invariance by imposing boundary conditions (bc's) on bulk fields involving a symmetry twist. If the twist includes an R-symmetry group action, then SUSY is softly broken by the SSSB mechanism [32, 33]. This SSSB is non-local from the higher-dimensional perspective, and is of an exceptionally soft type, similar to finite-temperature breaking of SUSY. In our case the twists will be maximal, ± 1 , and the underlying non-gravitational sector can be described as a 5D gauge theory compactified on a $S^1/(Z_2 \times Z'_2)$ orbifold. The 5th dimension, of physical length πR , is parameterised by $y \in [0, \pi]$, and branes sit at the inequivalent fixed points at $0, \pi R$.

Our 5D bulk theory is a SUSY theory containing the SM gauge fields, the first two families, and a pair of distinct Higgs hypermultiplets, H_u, H_d , (see Fig.1a). As the minimal SUSY in 5D corresponds to $N = 2$ 4D SUSY, the superpartners of these bulk states fill out $N = 2$ 4D multiplets, with each 5D vector implying both a 4D vector and chiral supermultiplet in the adjoint representation, $V_{5D}^a = \{V_{4D}^a, \bar{\Sigma}^a\}$ (with physical fields V_μ^a, λ^a and $\bar{\sigma}^a, \bar{\lambda}^a$) while the matter fields are hypermultiplets consisting of 4D chiral and anti-chiral multiplets $\Phi_{5D}^i = \{\phi^i, \bar{\phi}^i\}$ (physical fields φ_i, ψ_i and $\bar{\varphi}_i, \bar{\psi}_i$) [34–38].

The two Z_2 actions, at their respective fixed points at $0, \pi R$, break 5D SUSY to two *different* and incompatible $N = 1$ 4D SUSYs thus breaking SUSY completely in the 4D effective theory; the component field bc's are summarised in Table I. Due to the non-local nature of SSSB there are *no cutoff-dependent log enhancements* of the effective 4D soft terms. At $y = 0$ we localise the

	(+, +)	(+, -)	(-, +)	(-, -)
V_{5D}^a	V_μ^a	λ^a	$\bar{\lambda}^a$	$\bar{\sigma}^a$
$H_{u,d}$	$h_{u,d}$	$\psi_{h_{u,d}}$	$\bar{\psi}_{h_{u,d}}$	$\bar{h}_{u,d}$
$F_{i=1,2}$	ψ_{F_i}	φ_{F_i}	$\bar{\varphi}_{F_i}$	$\bar{\psi}_{F_i}$
$\Phi_{1,2}$	$\psi_{\Phi_{1,2}}$	$\varphi_{1,2}$	$\bar{\varphi}_{1,2}$	$\bar{\psi}_{\Phi_{1,2}}$

TABLE I: Bc's at $y = (0, \pi)$ for bulk fields of complete model with \pm corresponding to Neumann/Dirichlet. Only the $(+, +)$ fields have a zero mode, and the KK mass spectrum ($n \geq 0$) is: $m_n = n/R$ for $(+, +)$ fields; $(2n+1)/2R$ for $(+, -)$ and $(-, +)$; and $(n+1)/R$ for $(-, -)$. $\psi_{F_{1,2}}$ stands for all 1st/2nd generation fermions; φ_{F_i} their 4D $N = 1$ sfermion partners; barred states are the extra 5D $N = 1$ SUSY partners.

3rd generation fields. As the fixed points preserve only $N = 1$ 4D SUSY, these states are simply 4D chiral multiplets with no additional partners, and a localised Yukawa superpotential for up-like states is allowed

$$\delta(y)H_u(y)\left(\frac{\tilde{y}_t}{M_5^{1/2}}Q_3U_3^c + \frac{\tilde{y}_c}{M_5^{3/2}}Q_2(y)U_2^c(y) + \dots\right). \quad (1)$$

where \tilde{y}_i are dimensionless and the Yukawa couplings to bulk 1st/2nd generations are naturally suppressed compared to the brane-localized 3rd generation. We later return to the down-type Yukawas.

There is no need for a μ term linking $H_u H_d$ to lift the higgsinos. Instead, SSSB gives the higgsinos a large $1/2R$ mass by marrying ψ_{h_u} with $\bar{\psi}_{h_u}$. The SSSB bc's lift the Higgsinos while making *no contribution* to the scalar Higgs masses, avoiding the usual source of tree-level tuning.

After SSSB the brane-localised scalars pick up, at 1-loop, finite positive soft SUSY-breaking masses

$$\delta\tilde{m}_i^2 \simeq \frac{7\zeta(3)}{16\pi^4 R^2} \left(\sum_{I=1,2,3} C_I(i)g_I^2 + C_t(i)y_t^2 \right), \quad (2)$$

with $C(U_3) = \{4/9, 0, 4/3, 1\}$, $C(D_3) = \{1/9, 0, 4/3, 0\}$, $C(E_3) = \{1, 0, 0, 0\}$, $C(L_3) = \{1/4, 3/4, 0, 0\}$, $C(Q_3) = \{1/36, 3/4, 4/3, 1/2\}$, and for the Higgs bulk scalar zero mode $C(H_{u,d}) = \{1/4, 3/4, 0, 0\}$ [7].

In addition to the positive 1-loop EW contribution Eq.(2), the Higgs soft mass $\tilde{m}_{H_u}^2$ receives a comparable negative contribution at 2-loops from the $t\text{-}\tilde{t}$ sector. Ref. [12] performed a 2-loop 5D calculation of this term, and we have used RG methods to determine the leading 3-loop $\log(m_t R)$ -, $\log(m_{\tilde{t}_1}/m_t)$ -enhanced corrections, which are numerically important in determining the fate of EWSB [39]. As shown in Fig. 2, these minimal contributions do *not* so far lead to EWSB. Nevertheless, the model has attractive features: Compared to 4D theories the Higgs soft mass is more screened from SUSY-breaking as Eq.(2) involves a finite 1-loop factor with no log enhancement, SUSY breaking for all but the 3rd generation and Higgs scalar zero mode is *direct and universal*, and higgsinos are heavy without a large μ term.

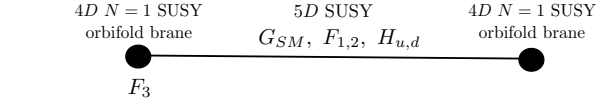


Fig. 1a

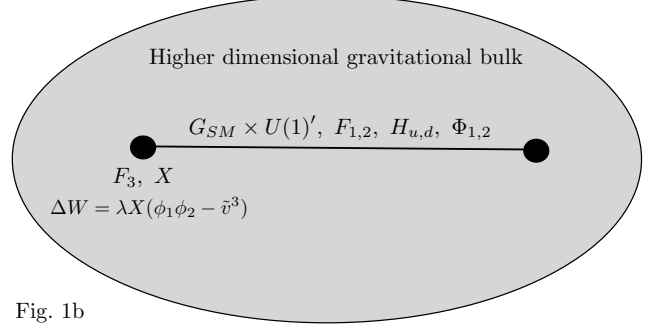


Fig. 1b

FIG. 1: (a) Schematic of minimal model. In 5D are the SM gauge fields, the first two families $F_{1,2}$, Higgs doublets $H_{u,d}$, and superpartners implied by 5D SUSY. The 3rd generation chiral multiplets are brane-localised. SUSY is broken non-locally by bc's. (b) Full model including embedding in yet higher-dimensional bulk. The 5D $U(1)'$ is broken via y -dependent VEVs (driven by the brane-localised superpotential ΔW) of bulk fields, $\Phi_{1,2}$, of charges ± 1 . After SSSB, $F_X \sim 1/R^2$ is induced for X , a brane-localised singlet field.

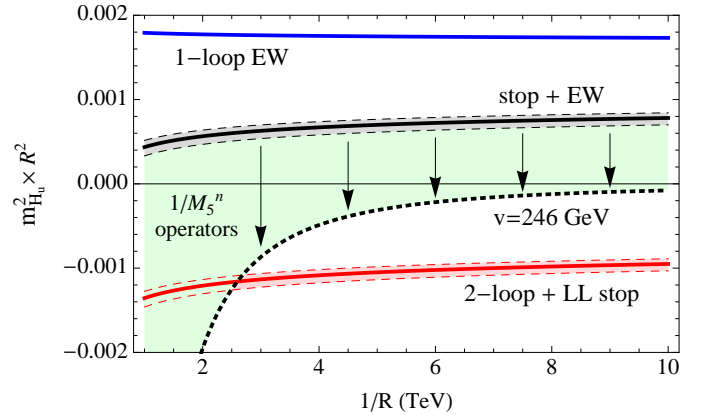


FIG. 2: Contributions to the Higgs soft mass $m_{H_u}^2$ in units of $1/R^2$. The positive 1-loop electroweak contribution (blue) and the negative 2-loop + leading log top-stop sector contribution (red) combine to give a positive mass squared (black). Contributions from higher-dimension operators Eq.(4) can lead to successful EWSB, indicated by the dotted black curve. The dashed bands show the uncertainty for $\overline{\text{MS}}$ top mass $m_t(M_t) = 160_{-4}^{+5}$ GeV.

SUCCESSFUL EWSB AND HIGGS MASS

Other faults remain in this model, and we find their solution plays a major role for EWSB and experimental signatures. First, our 5D theory is an effective theory which must be cutoff at a scale M_5 . The bulk 5D

gauge couplings are dimensionful ($1/g_{I,4}^2 = \pi R/g_{I,5}^2$ up to small brane-kinetic-term corrections), and 5D perturbative unitarity bounds for g_3 require $\pi M_5 R \lesssim 25$ [40, 41]. NDA strong coupling estimates for the brane-localized Yukawas give a similar bound [42].

This cutoff is large enough to justify the 5D viewpoint and the parameterization of UV effects in terms of higher dimensional operators, but the weakness of gravity in the low energy 4D theory, $M_{\text{Pl}} \gg M_5$, must still be explained. The two controllable possibilities of which we are aware are: (a) Embed the 5D theory in a 10 or 11D string/M-theory where some or all of the extra 5 or 6 *purely gravitational* dimensions are ‘large’, similar to the original brane-world proposal of Refs.[43–46] (see Fig.1b). Since our fundamental scale is $M_5 \gtrsim 30$ TeV, $n \geq 2$ extra dimensions is safe from cosmological, astrophysical, and laboratory constraints. (b) Utilise a little-string-theory construction with tiny string coupling [47].

Second, the radius R is unstabilized. Moreover, SSSB without radius stabilization is of no-scale type with zero CC at tree level [36, 48, 49], and generally radius stabilization yields a deep negative CC of order $\sim -1/R^4$ [50–54]. An additional positive SUSY breaking sector (which may or may not be sequestered from the radion at tree level) can tune the minimum to zero CC, and will generally involve a brane-localised field, X , with an F-term, $F_X \sim 1/R^2$ [39].

With this additional brane-localized SUSY breaking $F_X \neq 0$, the Kahler operators

$$\Delta \mathcal{K}_{m_H^2} = \delta(y) \frac{c_H}{M_5^3} X^\dagger X H_u^\dagger H_u \quad (3)$$

$$\Delta \mathcal{K}_{m_{\tilde{t}}^2} = \delta(y) X^\dagger X \left(\frac{c_Q}{M_5^2} Q_3^\dagger Q_3 + \frac{c_U}{M_5^2} U_3^{c\dagger} U_3^c \right) \quad (4)$$

can alter the H_u soft mass and trigger EWSB, either directly for $\Delta \mathcal{K}_{m_H^2}$ or radiatively through one-loop stop corrections for $\Delta \mathcal{K}_{m_{\tilde{t}}^2}$. When the 5D picture is under good control, $(\pi R M_5) \gtrsim 10$, the contribution from $\Delta \mathcal{K}_{m_{\tilde{t}}^2}$ dominates. As illustrated in Fig. 2, we find that for $F_X \sim 1/R^2$, and for $c_Q, c_U \sim 1$ this shift is sufficient to trigger EWSB at scales $1/2R \gtrsim 2$ TeV. The tuning involved will be seen to be exceptionally mild for present collider limits.

The bottom and tau Yukawas also result from F_X via the Kahler terms [31]

$$\delta(y) (H_u(y)^\dagger X^\dagger) \left(\frac{\tilde{y}_b}{M_5^{5/2}} Q_3 D_3^c + \dots \right). \quad (5)$$

The 1st and 2nd generation down-type Yukawas can be generated by similar higher dimensional Kahler operators or by superpotential couplings to h_u^\dagger on the $y = \pi$ brane [14]. Therefore H_d need not obtain a VEV, a dramatic simplification of the Higgs sector only possible because the cutoff scale is so low. Although H_d must be present to avoid a quadratically divergent Fayet-Iliopoulos (FI)

term [16, 55, 56], unlike the MSSM there is no need for there to be μ - or B_μ -terms that link H_u to H_d . The simplest option is to impose an unbroken Z_2 symmetry on H_d which forbids these unnecessary terms and eliminates potentially dangerous flavor-changing effects; H_d is a stable (neutral) particle in the spectrum in addition to the LSP, as in the inert doublet models [57–60].

For $m_{\tilde{t}_1} \gtrsim 3.5$ TeV, the radiative contributions to the physical Higgs mass may be large enough to accommodate $m_h = 126$ GeV [61]. For lighter stops, we obtain $m_h = 126$ GeV with a non-decoupling D-term (as only $\langle H_u \rangle \neq 0$, a NMSSM-like singlet interaction $S H_u H_d$ can not be employed as in ref. [62]). Specifically, we introduce a bulk $U(1)'$ gauging a subgroup of right-handed $SU(2)$ generated by T_{3R} under which H_u (and H_d) transform (the $U(1)'$ is anomaly-free if three light RHD neutrino superfields are introduced in the bulk; we find our theory allows a novel theory of neutrino mass generation [39]). To avoid suppression of the quartic, the breaking of the new gauge group must couple to large SUSY breaking F-terms [63–65]. It is natural to associate the breaking of the $U(1)'$ with the same dynamics that generates F_X , with the resulting Z' mass $\sim 1/R$.

A simple model where F_X is induced by SSSB and is associated with the breaking of the $U(1)'$ is obtained by introducing bulk hypermultiplet fields $\phi_{1,2}$ charged $\pm \frac{1}{2}$ under the $U(1)'$ with SSSB bc’s given in Table. 1 and a brane-localized superpotential

$$\Delta W = \frac{\lambda}{M_5} X (\phi_1(y) \phi_2(y) - \tilde{v}^3) \delta(y). \quad (6)$$

This leads to spontaneous breaking of the $U(1)'$ in the D-flat direction with a y -dependent profile for $\langle \phi_{1,2} \rangle$ and a brane-localized $F_X = M_5/(\lambda \pi R)$. This positive SUSY breaking contribution to the radion potential can be tuned to allow stabilization with zero CC. We find that for $m_{\tilde{t}} \gtrsim 650$ GeV ($m_{\tilde{t}} \gtrsim 1$ TeV) and $m'_{Z'} \lesssim 2/R$, the $U(1)'$ D-term can yield $m_h = 126$ GeV with $g_X < 1$ ($g_X < g_2$). The $U(1)'$ sector also contributes to the Higgs soft mass. The contribution is not well-approximated by the truncated lightest KK modes; we evaluate it in the 5D theory and find for $m_{Z'} \gtrsim 1/R$ the contribution favors EWSB and numerically approaches

$$\delta m_{H_u}^2 (U(1)') \approx -10^{-3} g_X^2 m_{Z'}^2. \quad (7)$$

PHENOMENOLOGY AND VARIATIONS

The theory has a rich phenomenology, and a variety of new physics signatures are accessible to LHC14 in the low-fine-tuning parameter region. Here we provide just a brief summary of the main features [39]. The spectrum of new (non-gravitational) states is illustrated in Fig. 3, where we have shown values with minimal fine-tuning consistent with current bounds.

for which the theory is $\sim 20\%$ tuned. For $m_{\tilde{t}_1} \gtrsim 3.5$ TeV, the tuning is still only at the few percent level and $m_h = 126$ GeV might be obtained radiatively [61] without the complications of an extra $U(1)'$ sector – an attractive target for a 100 TeV proton collider.

The production of KK excitations of SM particles would be an important signature of the extra-dimensional nature of this model, but their large mass $\sim 1/R$ and an approximate KK-parity make these particles difficult to reach at LHC14. Observing the near degeneracy of gauginos, higgsinos, and 1st/2nd generation sfermions would provide an alternative strong indication of the extra-dimensional nature of the theory.

In summary, we have presented a model where SSSB accompanied by a simple mechanism driving EWSB leads to a natural spectrum consistent with Higgs properties and sparticle bounds with fine-tuning better than $\sim 50\%$ even after LHC8 limits. Variations involving different field content and localizations, including interplay with other mechanisms for driving EWSB in SSSB via different bc's [10, 11] and quasi-localization of the stop [12–16] or Higgs [17–20] deserve further attention as leading candidates for natural theories at LHC14 and future colliders. In an aesthetic direction, the extended gauge structure and extra dimensions suggest interesting possibilities for gauge unification in this model [39].

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- [1] T. Gherghetta, B. von Harling, A. D. Medina, and M. A. Schmidt, JHEP **02**, 032 (2013), 1212.5243.
- [2] A. Arvanitaki, M. Baryakhtar, X. Huang, K. van Tilburg, and G. Villadoro, JHEP **1403**, 022 (2014), 1309.3568.
- [3] E. Hardy, JHEP **1310**, 133 (2013), 1306.1534.
- [4] J. L. Feng, Ann.Rev.Nucl.Part.Sci. **63**, 351 (2013), 1302.6587.
- [5] T. Gherghetta, B. von Harling, A. D. Medina, and M. A. Schmidt (2014), 1401.8291.
- [6] J. Fan and M. Reece (2014), 1401.7671.
- [7] I. Antoniadis, S. Dimopoulos, A. Pomarol, and M. Quiros, Nucl.Phys. **B544**, 503 (1999), hep-ph/9810410.
- [8] A. Delgado, A. Pomarol, and M. Quiros, Phys.Rev. **D60**, 095008 (1999), hep-ph/9812489.
- [9] A. Pomarol and M. Quiros, Phys.Lett. **B438**, 255 (1998), hep-ph/9806263.
- [10] A. Delgado and M. Quiros, Nucl.Phys. **B607**, 99 (2001), hep-ph/0103058.
- [11] A. Delgado, G. von Gersdorff, and M. Quiros, Nucl.Phys. **B613**, 49 (2001), hep-ph/0107233.
- [12] R. Barbieri, G. Marandella, and M. Papucci, Nucl.Phys. **B668**, 273 (2003), hep-ph/0305044.
- [13] R. Barbieri, L. J. Hall, G. Marandella, Y. Nomura, T. Okui, et al., Nucl.Phys. **B663**, 141 (2003), hep-ph/0208153.
- [14] R. Barbieri, G. Marandella, and M. Papucci, Phys.Rev. **D66**, 095003 (2002), hep-ph/0205280.
- [15] R. Barbieri, L. J. Hall, and Y. Nomura, Phys.Rev. **D63**, 105007 (2001), hep-ph/0011311.
- [16] D. Marti and A. Pomarol, Phys.Rev. **D66**, 125005 (2002), hep-ph/0205034.
- [17] D. Diego, G. von Gersdorff, and M. Quiros, JHEP **0511**, 008 (2005), hep-ph/0505244.
- [18] D. Diego, G. von Gersdorff, and M. Quiros, Phys.Rev. **D74**, 055004 (2006), hep-ph/0605024.
- [19] G. von Gersdorff, Mod.Phys.Lett. **A22**, 385 (2007), hep-ph/0701256.
- [20] G. Bhattacharyya and T. S. Ray, JHEP **1205**, 022 (2012), 1201.1131.
- [21] S. Dimopoulos and G. Giudice, Phys.Lett. **B357**, 573 (1995), hep-ph/9507282.
- [22] A. Pomarol and D. Tommasini, Nucl.Phys. **B466**, 3 (1996), hep-ph/9507462.
- [23] A. G. Cohen, D. Kaplan, and A. Nelson, Phys.Lett. **B388**, 588 (1996), hep-ph/9607394.
- [24] I. Antoniadis, S. Dimopoulos, and G. Dvali, Nucl.Phys. **B516**, 70 (1998), hep-ph/9710204.
- [25] R. Barbieri, L. J. Hall, and Y. Nomura, Nucl.Phys. **B624**, 63 (2002), hep-th/0107004.
- [26] A. Delgado, G. von Gersdorff, P. John, and M. Quiros, Phys.Lett. **B517**, 445 (2001), hep-ph/0104112.
- [27] R. Contino and L. Pilo, Phys.Lett. **B523**, 347 (2001), hep-ph/0104130.
- [28] N. Weiner (2001), hep-ph/0106021.
- [29] H. D. Kim, Phys.Rev. **D65**, 105021 (2002), hep-th/0109101.
- [30] M. Puchwein and Z. Kunszt, Annals Phys. **311**, 288 (2004), hep-th/0309069.
- [31] R. Davies, J. March-Russell, and M. McCullough, JHEP **1104**, 108 (2011), 1103.1647.
- [32] J. Scherk and J. H. Schwarz, Phys.Lett. **B82**, 60 (1979).
- [33] J. Scherk and J. H. Schwarz, Nucl.Phys. **B153**, 61 (1979).
- [34] N. Marcus, A. Sagnotti, and W. Siegel, Nucl.Phys. **B224**, 159 (1983).
- [35] N. Arkani-Hamed, T. Gregoire, and J. G. Wacker, JHEP **0203**, 055 (2002), hep-th/0101233.
- [36] D. Marti and A. Pomarol, Phys.Rev. **D64**, 105025 (2001), hep-th/0106256.
- [37] A. Hebecker, Nucl.Phys. **B632**, 101 (2002), hep-ph/0112230.
- [38] I. Lynch, William Divine, M. A. Luty, and J. Phillips, Phys.Rev. **D68**, 025008 (2003), hep-th/0209060.
- [39] S. Dimopoulos, K. Howe, and J. March-Russell (In preparation).
- [40] A. Muck, L. Nilse, A. Pilaftsis, and R. Ruckl, Phys.Rev. **D71**, 066004 (2005), hep-ph/0411258.
- [41] R. S. Chivukula, D. A. Dicus, H.-J. He, and S. Nandi, Phys.Lett. **B562**, 109 (2003), hep-ph/0302263.

- [42] G. Marandella and M. Papucci, Phys.Rev. **D71**, 055010 (2005), hep-ph/0407030.
- [43] N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys.Rev. **D59**, 086004 (1999), hep-ph/9807344.
- [44] I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys.Lett. **B436**, 257 (1998), hep-ph/9804398.
- [45] N. Arkani-Hamed, S. Dimopoulos, G. Dvali, and J. March-Russell, Phys.Rev. **D65**, 024032 (2002), hep-ph/9811448.
- [46] P. C. Argyres, S. Dimopoulos, and J. March-Russell, Phys.Lett. **B441**, 96 (1998), hep-th/9808138.
- [47] I. Antoniadis, A. Arvanitaki, S. Dimopoulos, and A. Givon, Phys.Rev.Lett. **108**, 081602 (2012), 1102.4043.
- [48] M. A. Luty and N. Okada, JHEP **0304**, 050 (2003), hep-ph/0209178.
- [49] D. E. Kaplan and N. Weiner (2001), hep-ph/0108001.
- [50] E. Ponton and E. Poppitz, JHEP **0106**, 019 (2001), hep-ph/0105021.
- [51] G. von Gersdorff, M. Quiros, and A. Riotto, Nucl.Phys. **B689**, 76 (2004), hep-th/0310190.
- [52] R. Rattazzi, C. A. Scrucca, and A. Strumia, Nucl.Phys. **B674**, 171 (2003), hep-th/0305184.
- [53] G. von Gersdorff and A. Hebecker, Nucl.Phys. **B720**, 211 (2005), hep-th/0504002.
- [54] E. Dudas and M. Quiros, Nucl.Phys. **B721**, 309 (2005), hep-th/0503157.
- [55] D. Ghilencea, S. Groot Nibbelink, and H. P. Nilles, Nucl.Phys. **B619**, 385 (2001), hep-th/0108184.
- [56] R. Barbieri, R. Contino, P. Creminelli, R. Rattazzi, and C. Scrucca, Phys.Rev. **D66**, 024025 (2002), hep-th/0203039.
- [57] M. Cirelli, N. Fornengo, and A. Strumia, Nucl.Phys. **B753**, 178 (2006), hep-ph/0512090.
- [58] R. Barbieri, L. J. Hall, and V. S. Rychkov, Phys.Rev. **D74**, 015007 (2006), hep-ph/0603188.
- [59] E. Ma, Phys.Rev. **D73**, 077301 (2006), hep-ph/0601225.
- [60] L. Lopez Honorez, E. Nezri, J. F. Oliver, and M. H. Tytgat, JCAP **0702**, 028 (2007), hep-ph/0612275.
- [61] J. L. Feng, P. Kant, S. Profumo, and D. Sanford, Phys.Rev.Lett. **111**, 131802 (2013), 1306.2318.
- [62] A. Delgado and M. Quiros, Phys.Lett. **B484**, 355 (2000), hep-ph/0004124.
- [63] P. Batra, A. Delgado, D. E. Kaplan, and T. M. Tait, JHEP **0402**, 043 (2004), hep-ph/0309149.
- [64] A. Maloney, A. Pierce, and J. G. Wacker, JHEP **0606**, 034 (2006), hep-ph/0409127.
- [65] C. Cheung and H. L. Roberts, JHEP **1312**, 018 (2013), 1207.0234.
- [66] S. Dimopoulos, I. Garcia Garcia, K. Howe, and J. March-Russell (In preparation).
- [67] P. Ade et al. (BICEP2 Collaboration) (2014), 1403.3985.
- [68] N. Arkani-Hamed, S. Dimopoulos, N. Kaloper, and J. March-Russell, Nucl.Phys. **B567**, 189 (2000), hep-ph/9903224.
- [69] C. Cheung, F. D'Eramo, and J. Thaler, JHEP **1108**, 115 (2011), 1104.2600.
- [70] G. Aad et al. (ATLAS), JHEP **1310**, 189 (2013), 1308.2631.
- [71] S. Chatrchyan et al. (CMS Collaboration), Eur.Phys.J. **C73**, 2677 (2013), 1308.1586.
- [72] Tech. Rep. ATLAS-CONF-2013-017, CERN, Geneva (2013).
- [73] S. Chatrchyan et al. (CMS Collaboration), Phys.Lett. **B720**, 63 (2013), 1212.6175.
- [74] Tech. Rep. ATL-PHYS-PUB-2013-011, CERN, Geneva (2013).
- [75] Tech. Rep. ATLAS-CONF-2013-061, CERN, Geneva (2013).
- [76] S. Chatrchyan et al. (CMS Collaboration) (2013), 1311.4937.